Vacuum phase transition at nonzero baryon density.

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March 27, 2007

Abstract

It is argued that the dominant contribution to the interaction of quark gluon plasma at moderate $T \geq T_c$ is given by the nonperturbative vacuum field correlators. Basing on that nonperturbative equation of state of quark-gluon plasma is computed and in the lowest approximation expressed in terms of absolute values of Polyakov lines for quarks and gluons $L_{fund}(T); L_{adj}(T) = (L_{fund})^{9/4}$known from lattice and analytic calculations. Phase transition at any μ is described as a transition due to vanishing of one of correlators, $D^E(x)$, which implies the change of gluonic condensate ΔG_2. Resulting transition temperature $T_c(\mu)$ is calculated in terms of ΔG_2 and $L_{fund}(T_c)$.

The phase curve $T_c(\mu)$ is in good agreement with lattice data. In particular $T_c(0) = 0.27; 0.19; 0.17$ GeV for $n_f = 0, 2, 3$ and fixed $\Delta G_2 = 0.0035$ GeV^4.

1 Introduction

QCD at nonzero temperature T and quark chemical potential μ, and specifically phase transitions are important topics both for theory, and experiment [1, 2] and are carefully studied on the lattice (for recent reviews see [3, 4]).

Since these phenomena cannot be explained in perturbative QCD, any analytic approach should be based on nonperturbative methods, and as such solvable models, like NJL, have been used mostly. However the models used do not contain confinement and treat phase transition from the chiral symmetry (CS) point of view. Therefore it is interesting to look at EoS and phase transition in the framework of a nonperturbative (NP) method, based on vacuum fields and implementing confinement.

In this letter we are using such tool, which is called the Field Correlator Method (FCM) [5]. In FCM all NP information about QCD vacuum is contained in field correlators, and the Casimir scaling, now proved on the lattice also for $T \geq T_c$ [6], allows (with $\sim 1\%$ accuracy) to consider only four quadratic correlators for $T, \mu > 0$ of colorelectric (colormagnetic) field strengths, $D^E(x), D_1^E(x), (D^H(x), D_1^H(x))$. The correlators were calculated on the lattice [7, 8] and also analytically [9].

Most calculations with FCM at $T = \mu = 0$ contain as input only string tension σ, strong coupling $\alpha_s(q)$ and current (pole) quark masses, and results for hadronic spectra are in a good agreement with available experimental and lattice data, see last ref. in [5] and [10] for reviews.

In [11] the FCM was extended to the case of nonzero T, and in [12] the phase transition at $\mu = 0$ was described as a Vacuum Dominated (VD) transition from the confining vacuum (all $D^{(i)}, D_1^{(i)}, i = E, H$ are nonzero, and $D^E = D^H, D_1^E = D_!^H$ as for $T = 0$) to the deconfined state (only D^E and hence $\sigma^E = \frac{1}{2}\int D^E(x)d^2x$ are zero). This picture was supported by lattice measurements [7, 8] where indeed D^E was shown to vanish at $T = T_c$, while D_1^E and colormagnetic correlators stay intact. This fact supports the conjecture [13] that D_1^E is nonzero at $T > T_c$ and strong enough to support bound states of quarks and gluons. Bound states of $q\bar{q}$ were indeed found on the lattice years later (see [4] for a review) and in [14, 15] the connection of those to the correlator D_1^E was quantitatively established.

Moreover in [14] it was shown, that a new static potential $V_1(r, T)$ is provided by D_1^E at $T \geq T_c$ and the NP part of it yields the modulus of the renormalized Polyakov line, $L_{fund} = exp\left(-\frac{V_1(\infty,T)}{2T}\right)$.

The physical picture of the VD deconfining phase transition given in [12], has allowed to get the (simplified) estimate of the transition temperature T_c at $\mu = 0$ in reasonable agreement with lattice data for $n_f = 0, 2, 3$ the only input being the standard value of the gluonic condensate G_2, known from the ITEP sum rule method [16]. It is the purpose of the present letter to extend the method to the case of nonzero μ and nonzero interaction with vacuum ($V_1 \neq 0$) and to find $T_c(\mu)$, for $\mu > 0$.

To this end we are using the EoS found recently in [17], where pressure in the lowest approximation is expressed through the only dynamical quantity – $L_{fund}(T)$ calculated via the correlator $D_1^E(x)$ or obtained from lattice data. This factor takes into account interaction of the quark with the NP vacuum in the form of $L_{fund}(T)$ and analogous factor for gluons is $L_{adj}(T) = (L_{fund})^{9/4}$. The interaction between quarks and gluons is considered as a next step, and is argued to possibly contribute in a narrow temperature region near T_c [17, 18]. This simple picture of VD dynamics with the only input $L_{fund}(T)$, taken from lattice or analytic calculation suggested in [17], which may be called Vacuum Dominance Model (VDM), is adopted below and is shown to produce surprisingly reasonable results, being in good agreement with available lattice data for $\mu = 0$ or nonzero μ, where these data are reliable.

2 Nonperturbative EoS for $\mu > 0$

The main idea of the VDM discussed below is that the most important part of quark and gluon dynamics in the *sqgp* is the interaction of each individual quark or gluon with vacuum fields. This interaction is derived from field correlators and is rigorously proved to be embodied in factors, which happen to coincide with the modulus of Polyakov loop[1],

$$L_{fund} = \exp(-\frac{V_1(T) + 2V_D}{2T}), \quad L_{adj} = \exp(-\frac{9}{8}\frac{V_1(T) + 2V_D}{2T}) \tag{1}$$

where $V_1(T) \equiv V_1(\infty, T)$, $V_D \equiv V_D(r^*, T)$ and $V_1(r, T), V_D$ found in [14] to be $(\beta \equiv 1/T)$

$$V_1(r, T) = \int_0^\beta d\nu(1 - \nu T) \int_0^r \xi d\xi D_1^E(\sqrt{\xi^2 + \nu^2}), \tag{2}$$

$$V_D(r, T) = 2\int_0^\beta d\nu(1 - \nu T) \int_0^r (r - \xi) d\xi D^E(\sqrt{\xi^2 + \nu^2}). \tag{3}$$

In (1) r^* is the m.s.r. of the heavy-light $Q\bar{q}$ or its adjoint equivalent system; for $T > T_c$ one has $D^E = V_D = 0$. For $T < T_c$ one has $r^*_{fund} = \infty$

[1] We neglect in the approximation the difference between L_{fund} expressed via $V_1(\infty, T)$ and L^{lat}_{fund} where the role of $V_1(T)$ is played [14, 17] by singlet $Q\bar{Q}$ free energy $F^1_{Q\bar{Q}}(\infty, T)$. The latter quantity contains all excited states, so that $V_1(T) \geq F^1_{Q\bar{Q}}(\infty, T)$.

for $n_f = 0$, yielding $L_{fund} = 0$, however $r^*_{adj} \approx 0.4$ fm for any n_f and gives nonzero $L_{adj}(T < T_c)$. The form (1)-(3) is in good agreement with lattice data [6] and also explains why L_{fund} is a good order parameter (however approximate for $n_f \neq 0$). Note, that only NP parts D_1^E, D^E enter in (2), (3), see first ref. in [9] for discussion of separation of these parts; renormalization procedure is discussed in [14, 15].

In the lowest NP approximation one neglects pair, triple etc. interactions between quarks and gluons (which are important for $T_c \leq T \leq 1.2T_c$ where density is low and screening by medium is not yet operating, see [17, 18])) and derives the following EoS (this approximation is called in [17] the Single Line Approximation (SLA))

$$p_q \equiv \frac{P_q^{SLA}}{T^4} = \frac{4N_c n_f}{\pi^2} \sum_{n=1}^{\infty} \frac{(-1)^{n+1}}{n^4} L^n_{fund} \varphi_q^{(n)} cosh \frac{\mu n}{T} \tag{4}$$

$$p_{gl} \equiv \frac{P_{gl}^{SLA}}{T^4} = \frac{2(N_c^2-1)}{\pi^2} \sum_{n=1}^{\infty} \frac{1}{n^4} L^n_{adj} \tag{5}$$

with

$$\varphi_q^{(n)}(T) = \frac{n^2 m_q^2}{2T^2} K_2\left(\frac{m_q n}{T}\right) \approx 1 - \frac{1}{4}\left(\frac{n m_q}{T}\right)^2 + ... \tag{6}$$

In (4), (5) it was assumed that $T \lesssim \frac{1}{\lambda} \cong 1$ GeV, where λ is the vacuum correlation length, e.g. $D_1^{(E)}(x) \sim e^{-|x|/\lambda}$, hence powers of L_i^n, see [17] for details.

With few percent accuracy one can replace the sum in (5) by the first term, $n = 1$, and this form will be used below for p_{gl}, while for p_q this replacement is not valid for large $\frac{\mu}{T}$, and one can use instead the form equivalent to (4),

$$p_q = \frac{n_f}{\pi^2}\left[\Phi_\nu\left(\frac{\mu - \frac{V_1}{2}}{T}\right) + \Phi_\nu\left(-\frac{\mu + \frac{V_1}{2}}{T}\right)\right] \tag{7}$$

where $\nu = m_q/T$ and

$$\Phi_\nu(a) = \int_0^\infty \frac{z^4 dz}{\sqrt{z^2+\nu^2}} \frac{1}{(e^{\sqrt{z^2+\nu^2}-a}+1)}. \tag{8}$$

Eqs. (7), (5) define p_q, p_{gl} for all T, μ and m_q, which is the current (pole) quark mass at the scale of the order of T.

To draw p_q, p_{gl} and $p \equiv p_q + p_{gl}$ as functions of T, μ one needs explicit form of $V_1(T)$. This was obtained analytically and discussed in [14]; another form was found from $D_1^E(x)$ measured on the lattice in [8, 9] and is given in [15].

Also from lattice correlator studies [8, 9], [14] $V_1(T = T_c)$ is (with $\sim 10\%$ accuracy) 0.5 GeV and is decreasing with the growth of T (cf. Fig. 2 of [14] and Fig. 1 of [15]). This behaviour is similar to that found repeatedly on the lattice direct measurement of $F_\infty^{(1)}$, see e.g. Fig.2 of [19] where $F_\infty^{(1)} = V_1(T)$ is given for $n_f = 0, 2, 3$. In what follows we shall exploit the latter curves parametrizing them for $T \geq T_c$ and all n_f as

$$V_1(T) = \frac{0.175 \text{ GeV}}{1.35\left(\frac{T}{T_c}\right) - 1}, \quad V_1(T_c) \approx 0.5 \text{ GeV}. \tag{9}$$

For $\mu > 0$ one can expect a μ-dependence of V_1, however it should be weak for values of μ much smaller than the scale of change of vacuum fields. The latter scale can be identified with the dilaton mass m_d, which is of the order of the lowest glueball mass, i.e. ~ 1.5 GeV $\equiv m_d$. Hence one can expect, that V_1 in the lowest approximation does not depend on μ. This is supported by the lattice measurements in [20], where for $\frac{T}{T_c} = 1.5$ and $\frac{\mu}{T} = 0.8$ the values of $F_\infty^{(1)}$ are almost indistinguishable from the case of $\frac{\mu}{T} = 0$.

To give an illustration of the resulting EoS we draw in Fig.1 the pressure p for the cases $\mu = 0$, $n_f = 0, 2$. One can see a reasonable behaviour for $T_c(n_f) = (0.27; 0.19; 0.17)$ GeV (for $n_f = 0, 2, 3$ respectively)similar to the lattice data, see [21] for a review and discussion.

3 Phase transition for nonzero μ

Here we extend the VD mechanism suggested in [12] to the case of nonzero μ and V_1. We assume as was said in Introduction that the phase transition occurs from the full confining vacuum with all correlators D^E, D_1^E, D^H, D_1^H present to the deconfined vacuum where D^E vanishes. The basics of our physical picture is that all fields (and correlators) do not change for μ, T changing in a wide interval, unless μ, T become comparable to the dilaton mass $m_d \approx M(\text{glueball } 0^+) \approx 1.5$ GeV[2]. Therefore correlators and

[2]Note that G_2 does not depend on μ, n_f in the leading order of the $1/N_c$ expansion and one expects a growth of the magnetic part of G_2 at $T > T_{dim.red} \approx 2T_c$, $G_2^{mag} \approx O(T^4)$,

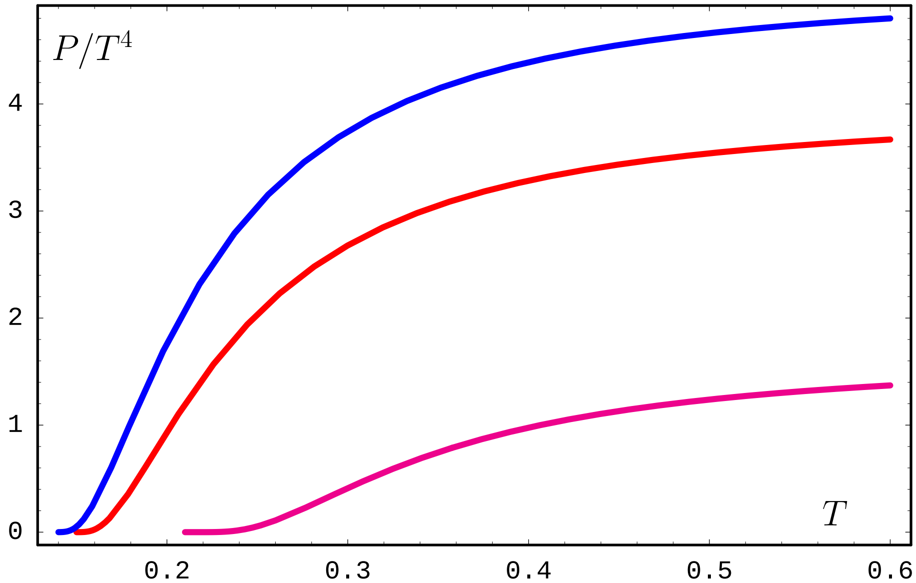

Figure 1: Pressure $\frac{P}{T^4}$ from Eq.(4,5) as function of temperature T (in GeV) for $n_f = 3, 2, 0$ (top to bottom) and $\Delta G_2 = 0.0034$ GeV^4.

σ^E, σ^H are almost constant till $T = T_c$ and at $T \geq T_c$ a new vacuum phase with $D^E = \sigma^E = 0$ is realised,which yields lower thermodynamic potential (higher pressure). Lattice measurements [7, 8] support this picture. The crucial step is that one should take into account in the free energy F of the system also the free energy of the vacuum, i.e. vacuum energy density $\varepsilon_{vac} = \frac{1}{4}\theta_{\mu\mu} = \frac{\beta(\alpha_s)}{16\alpha_s}\langle (F^a_{\mu\nu})^2\rangle = -\frac{(11-\frac{2}{3}n_f)}{32} G_2$, which can be estimated via the standard gluonic condensate [16] $G_2 \equiv \frac{\alpha_s}{\pi}\langle (F^a_{\mu\nu})^2\rangle \approx 0.012$ GeV^4.

Hence for the pressure $P = -F$ one can write in the phase I (confined)

$$P_I = |\varepsilon_{vac}| + \chi_1(T) \tag{10}$$

where $\chi(T)$ is the hadronic gas pressure, starting with pions, $\chi_{pion} \cong \frac{\pi^2}{30} T^4$. In the deconfined phase one can write

$$P_{II} = |\varepsilon^{dec}_{vac}| + (p_{gl} + p_q) T^4 \tag{11}$$

where $|\varepsilon^{dec}_{vac}|$ is the vacuum energy density in the deconfined phase, which is mostly (apart from $D^E_1(0) \approx 0.2 D^E(0)$ see [7] and first ref. of [9]) colormag-

in the regime of dimensional reduction.

netic energy density and by the same reasoning as before we take it as for $T=0$, i.e. $|\varepsilon_{vac}^{dec}| \cong 0.5|\varepsilon_{vac}|$.

Equalizing P_I and P_{II} at $T=T_c(\mu)$ one obtains the equation for T_c

$$T_c(\mu) = \left(\frac{\Delta|\varepsilon_{vac}| + \chi(T)}{p_{gl}+p_q}\right)^{1/4} \tag{12}$$

where $\Delta|\varepsilon_{vac}| = |\varepsilon_{vac}| - |\varepsilon_{vac}^{dec}| \approx \frac{(11-\frac{2}{3}n_f)}{32}\Delta G_2$; $\Delta G_2 \approx \frac{1}{2}G_2$; p_{gl} and p_q are given in (5), (7) respectively and depend on both T_c and μ.

In this letter we shall consider the simplest case when the contribution of hadronic gas $\chi_1(T)$ can be neglected in the first approximation. Indeed, pionic gas yields only $\sim 7\%$ correction to the numerator of (12) at $T\approx T_c$, and from [22] one concludes that $\chi(T_c) \lesssim 0.5T_c^4$, which yields a $\lesssim 10\%$ increase of T_c for $G_2 \sim 0.01$ GeV.

From the expression for $T_c(\mu)$ (12) one can find limiting behaviour of $T_c(\mu\to 0)$ and $\mu_c(T\to 0)$. For the first one can use for p_q and p_g (4) and (5) and expand r.h.s. of (12) in ratio p_g/p_q with the result.

$$T_c = T^{(0)}\left(1+\frac{V_1(T_c)}{8T_c} + O\left(\left(\frac{V_1(T_c)}{8T_c}\right)^2\right)\right) \tag{13}$$

where the last term yields a 3% correction, and $T^{(0)} = \left(\frac{(11-\frac{2}{3}n_f)\pi^2\Delta G_2}{32\cdot 12 n_f}\right)^{1/4}$. Solving (13) for T_c one has

$$T_c = \frac{1}{2}T^{(0)}\left(1+\sqrt{1+\frac{\kappa}{T^{(0)}}}\right)\left(1+\frac{m_q^2}{16T_c^2}\right) \tag{14}$$

with $\kappa \equiv \frac{1}{2}V_1(T_c)$. From (12), (13) one can compute expansion $T_c(\mu)$ in powers of μ,

$$T_c(\mu_B) = T_c(0)\left(1 - C\frac{\mu_B^2}{T_c^2(0)}\right), \quad \mu_B = 3\mu.$$

$$C = \frac{1+\sqrt{1+\frac{\kappa}{T^{(0)}}}}{144\sqrt{1+\frac{\kappa}{T^{(0)}}}} = 0.0110(3) \quad \text{for } n_f = 2,3,4.$$

One can see, that C practically does not depend on n_f and is in the same ballpark as the values found by lattice calculations, see [23], [3] for reviews and references.

Another end point of the phase curve, $\mu_c(T \to 0)$, is found from (12) when one takes into account asymptotics $\Phi_0(a \to \infty) = \frac{a^4}{4} + \frac{\pi^2}{2}a^2 + \frac{7\pi^4}{60} + ...$, which yields (for small $\frac{m_q}{\mu}$)

$$\mu_c(T \to 0) = \frac{V_1(T_c)}{2} + (48)^{1/4} T^{(0)} \left(1 + \frac{3m_q^2}{4\mu_c^2}\right) \left(1 - \frac{\pi^2}{2} \frac{T^2}{\left(\mu_c - \frac{V_1(T_c)}{2}\right)^2}\right) \tag{15}$$

The resulting curve $T_c(\mu)$ according to Eq.(12) with $\chi_1 \equiv 0$ is given in Fig. 2 for$\Delta G_2 = 0.00341 \mathrm{GeV}^4$, $n_f = 2, 3$ and $m_q = 0$.

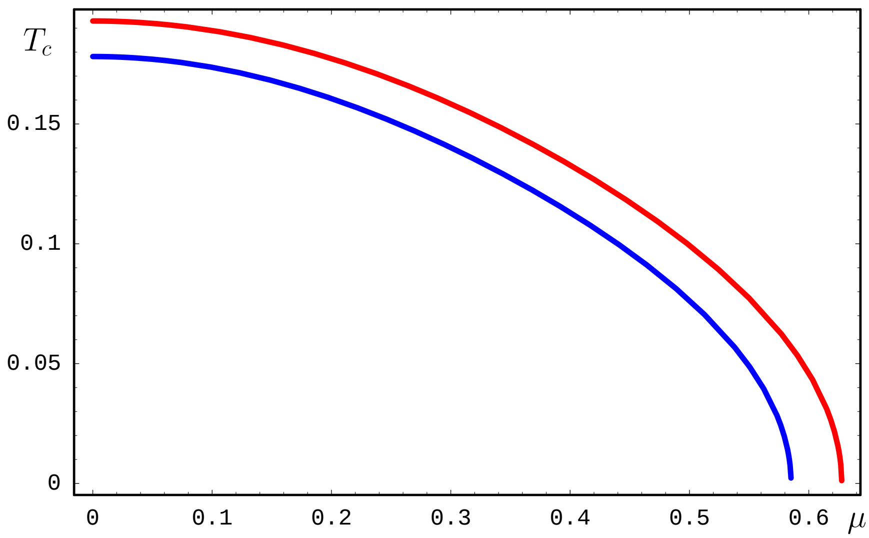

Figure 2: The phase transition curve $T_c(\mu)$ from Eq.(12) (in GeV) as function of quark chemical potential μ (in GeV) for $n_f = 2$ (upper curve) and $n_f = 3$ (lower curve) and $\Delta G_2 = 0.0034$ GeV^4.

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4 Discussion of results

Our prediction of $T_c(\mu)$ depends only on two numbers: 1) the value of gluonic condensate ΔG_2 and 2) the value of $V_1(T_c) = 0.5$ GeV taken from lattice data [19] (and quantitatively close to the value from the analytic form [14]).

We take G_2 in the limits $0.004\ \mathrm{GeV}^4 \le G_2 \le 0.015\ \mathrm{GeV}^4$, the value $G_2 = 0.008\mathrm{GeV}^4$ ($\Delta G_2 = 0.0034\ \mathrm{GeV}^4$) being in agreement with lattice data of $T_c(0)$, for $n_f = 0, 2, 3$, see Table 1.

Table 1. The values of $T_c(\mu = 0)$ and $\mu_c(T = 0)$ computed using (14) and (15) for several values of ΔG_2 and $n_f = 0, 2, 3$

$\frac{\Delta G_2}{0.01\ \mathrm{GeV}^4}$	0.191	0.341	0.57	1
T_c(GeV) $n_f = 0$	0.246	0.273	0.298	0.328
T_c(GeV) $n_f = 2$	0.168	0.19	0.21	0.236
T_c(GeV) $n_f = 3$	0.154	0.172	0.191	0.214
μ_c(GeV) $n_f = 2$	0.576	0.626	0.68	0.742
μ_c(GeV) $n_f = 3$	0.539	0.581	0.629	0.686

Note that T_c at $n_f = 0$ in Table 1 is obtained not from (14), but directly from (12) with $n_f = 0, \chi(T) = 0$.

The curve in Fig.2 has the expected form, which agrees with the curve, obtained in [24] by the reweighting technic and agrees for $\mu < 300$ MeV with that, obtained by the density of state method [25], and by the imaginary μ method [26].

An analysis of the integral (8) for $\nu = 0$ reveals that it has a mild singular point at $\mu_{sing} = \frac{V_1}{2} \pm i\pi T$, which may show up in derivatives in $\frac{\mu}{T}$. At $T = 0, \mu_{sing} = \frac{V_1}{2} \cong 0.25$ GeV and is close to the point where one expects irregularities on the phase diagram [25, 27].

The limit of small μ is given in (14). Taking $V_1(T_c) \approx 0.5$ GeV as follows from lattice and analytic estimates, one obtains with $\sim 3\%$ accuracy the values of T_c given on Fig.2 for $n_f = 2, 3$ and $\Delta G_2 = 0.0034$ GeV^4.

The values of μ_c, from (15) are given in Table 1 and are in agreement with the curves shown in Fig.3. Note that $\chi(T = 0) = 0$ and (15) holds also in the case, when hadron (and baryon) gas is taken into account.

At this point one should stress that our calculation in VDM of $p(T)$ does not contain model parameters and the only approximation is the neglect of interparticle interaction as compared to the interaction of each one with the vacuum (apart from neglect of $\chi(T)$). Fig.1 demonstrates that this approximation is reasonably good and one expects some 10÷15% accuracy in prediction of $T_c(\mu)$. Note that in VDM the phase transition is of the first order, which is supported for $n_f = 0$ by lattice data, see e.g. [6] however for $n_f = 2, 3$ lattice results disagree, see [28, 29] for a possible preference of the first order transition. One however should have in mind that the final conclusion for lattice data depends on input quark masses and continuum limit.

The "weakening" of the phase transition for $n_f > 0$ and nonzero quark masses is explained in our approach by the flattening of the curve $P(T)$ at $T \approx T_c$ when hadronic gas $\chi(T)$ is taken into account, since $\chi(T)$ for $n_f = 0$ is much smaller than for $n_f = 2$.

A few words about chiral symmetry properties of the transition. It was argued in a series of papers of one of the authors, that Chiral Symmetry Breaking (CSB) and confinement are closely connected, and moreover the effective scalar quark mass operator was calculated in terms of $D^E(x)(\sigma^E)$ in [30] and chiral condensate $\langle \bar{q}q \rangle$ and f_π via σ^E in [31].

Therefore $\langle \bar{q}q \rangle$ and f_π disappear at the same T_c where $D^E(x)$ (and σ^E) vanishes. This fact is in perfect agreement with known lattice data and supports the results of the present study.

Finally, as seen from our expressions for p_q, p_g (4), (7), our EoS is independent of $Z(N_c)$ factors and $Z(N_c)$ symmetry is irrelevant for the NP dynamics in our approach. This result is a consequence of more general property – the gauge invariance of the partition function for all T, μ, which requires that only closed Wilson loops appear in the resulting expressions, yielding finally only absolute value of the Polyakov loop, or in other words, is expressed via only singlet free energy of quark and antiquark $F^{(1)}_{Q\bar{Q}} \approx V_1$.

5 Summary and conclusions

We have calculated the QCD phase diagram in the lowest order of the background perturbation theory taking into account the change of gluonic vacuum condensate in the phase transition and modification of individual quark and gluon propagators in vacuum fields. Some support of this independent single particle approach can be found in the lack of S, B and S, Q correlations on the lattice [32] for $T \geq 1.1T_c$.

The resulting phase curve $T_c(\mu)$ depends only on two fundamental parameters: 1) the change of gluonic condensate across phase boundary, $\Delta G_2 \approx \frac{1}{2}G_2$ and 2) the absolute value of Polyakov loop at $T = T_c, L_{fund}(T_c) = \exp\left(-\frac{V_1(T_c)}{2T_c}\right)$, or else singlet $Q\bar{q}$ free energy $V_1(T_c) \cong F^1_{Q\bar{Q}}(\infty, T)$.

The second quantity is known from lattice and analytic calculations, $V_1(T_c) \approx 0.50(5)$ GeV.

It is shown that $T_c(\mu)$ depends on ΔG_2 rather mildly and yields reasonable values of $T_c(0) \approx 0.2$ GeV for generally accepted $G_2 \approx 0.01$ GeV^4 with $\mu_c(T = 0) \approx 0.6$ GeV. Moreover for the same value of $\Delta G_2 = 0.00341$ GeV^4 one obtains the set of three values $T_c(0) = (0.27; 0.19; 0.17)$ GeV^4 all being in agreement with lattice data.

Our results do not contain model or fitting parameters and allow in this (however crude) approximation connect in a very simple way the phase boundary to the fundamental properties of the QCD vacuum.

Many additional quantities like quark number susceptibilities, baryon - strangeness correlation etc. can be easily calculated in the method and shall be given elsewhere, as well as influence of possible diquark pairing.

The authors are grateful for discussions to N.O.Agasian, K.G.Boreskov, B.L.Ioffe, A.B.Kaidalov, O.V.Kancheli, B.O.Kerbikov.

The financial support of RFBR grants 06-02-17012, grant for scientific schools NSh-843.2006.2 and the State Contract 02.445.11.7424 is gratefully acknowledged. This work was done with financial support of the Federal Agency of Atomic Energy.

References

[1] J.Kapusta, Finite Temperature Field theory, Cambridge University Press, Cambridge, 1989; K.Yagi, T.Hatsuda and Y.Miake, Quark-gluon Plasma, Cambridge, Univ. Press, Cambridge 2005.

[2] B.Müller, J.L.Nagle, Ann. Rev. Nucl. Part. Phys. 1, (2006); nucl-th/0602029; D.d'Enterria, nucl-ex/0611012; BRAHMS: nucl-ex/0410020; PHENIX: nucl-ex/0410003; PHOBOS: nucl-ex/0410022; STAR: nucl-ex/0501009.

[3] C.Schmidt, hep-lat/0701019; C.Bernard et al., hep-lat/0611031; R.Gavai, hep-ph/0607050; U.M.Heller, hep-lat/0610114; F.Karsch, hep-lat/0601013, hep-lat/0608003.

[4] P.Petreczky, hep-lat/0409139; hep-lat/0609040.

[5] H.G.Dosch, Phys. Lett **B190** (1987), 177; H.G.Dosch, Yu.A.Simonov, Phys. Lett. **B 205** (1988), 339; Yu.A.Simonov, Nucl. Phys. **B 307** (1988), 512; A.Di Giacomo, H.G.Dosch, V.I.Shevchenko, Yu.A.Simonov, Phys. Rep. **372** (2002), 319.

[6] S.Gupta, K.Húbner, O.Kaczmarek, hep-lat/0608014.

[7] A.Di Giacomo and H.Panagopoulos, Phys. Lett. **B 285**(1992), 133 ;
M.D'Elia, A.Di Giacomo and E.Meggiolaro, Phys. Lett. **B 408** (1997), 315 ;
A.Di Giacomo, E.Meggiolaro and H.Panagopoulos, Nucl. Phys. **B 483**(1997), 371;
G.S.Bali, N.Brambilla, A.Vairo, Phys. Lett. **B 421** (1998), 265.

[8] M.D'Elia, A.Di Giacomo and E.Meggiolaro, Phys. Rev. **D 67** (2003), 114504.

[9] Yu.A.Simonov, Phys. At. Nucl. **69** (2006), 528 hep-ph/0501182; H.G.Dosch, M.Eidemúller, M. Jamin, Phys. Lett **B 452** (1999), 379; Yu.A.Simonov, Phys. At. Nucl. **61** (1998), 855, hep-ph/9712250.

[10] A.M.Badalian, V.I.Shevchenko, Yu.A.Simonov, Phys. At. Nucl. **69** (2006), 1781.

[11] Yu.A.Simonov, Phys. At. Nucl. **58** (1995), 309; hep-ph/9311216; Yu.A.Simonov, In "Varenna 1995, Selected topics in nonperturbative QCD", p.319; hep-ph/9509404; E.L.Gubankova and Yu.A.Simonov, Phys. Lett. **360** (1995), 93.

[12] Yu.A.Simonov, JETP Lett. **55** (1992), 627.

[13] Yu.A.Simonov, JETP Lett. **54** (1991), 249.

[14] Yu.A.Simonov, Phys. Lett. **B 619** (2005), 293.

[15] A.Di Giacomo, E. Meggiolaro, Yu.A.Simonov, A.I.Veselov, hep-ph/0512125, Phys.At. Nucl. (in press).

[16] M.Shifman, A.Vainshtein, V.Zakharov, Nucl. Phys. **B 147** (1979), 385, 448.

[17] Yu.A.Simonov, hep-ph/0702266.

[18] A.V.Nefediev, Yu.A.Simonov, in preparation

[19] O.Kaczmarek, F.Zantov, hep-lat/0506019.

[20] M.Döring, S.Ejiri, O.Kaczmarek, F.Karsch, E.Laermann, hep-lat/0509001.

[21] F.Karsch, J. Phys. Conf. Ser. **46** (2006), 122, hep- lat/0608003.

[22] F.Karsch, Prog. Theor. Phys. Suppl. **153** (2004), 106.

[23] S.Aoki, hep-lat/0509068.

[24] C.R.Allton, S.Ejiri, S.J.Hands, et al., hep-lat/0504011.

[25] Z.Fodor, S.D.Katz, C.Schmidt, hep-lat/0701022.

[26] M.P.Lombardo, hep-lat/0612017; M.D.'Elia and M.P.Lombardo, Phys. Rev. **D67** (2003), 014505, ibid **D70** (2004), 074509; Ph. de Forcrand and O.Philipsen, Nucl. Phys. **B673** (2003), 170.

[27] M.Stephanov, hep-lat/0701002.

[28] A.Di Giacomo, hep-lat/0610027.

[29] V.G.Bornyakov et al, hep-lat/0509122.

[30] Yu.A.Simonov, Phys. Atom. Nucl., **60** (1997), 2069, hep-ph/9704301; Phys. Rev. **D65** (2002), 094018.

[31] Yu.A.Simonov, Phys. Atom. Nucl. **67** (2004), 1027, hep-ph/0305281.

[32] R.V.Gavai, S.Gupta, Phys. Rev. **D73** (2006), 014004.